Precise measurement of $\Gamma(H \longrightarrow \gamma\gamma)$ at a PLC and theoretical consequences

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Abstract

With the LEP II Higgs search approaching exclusion limits on low values of $\tan\beta \sim 2$ it becomes increasingly important to investigate physical quantities sensitive to large masses of a pseudoscalar Higgs mass. This regime is difficult and over a large range of $\tan\beta$ impossible to cover at the LHC proton proton collider. In this paper we focus on the achievable statistical precision of the Higgs decay into two photons at a future $\gamma\gamma$ collider (PLC) in the MSSM mass range below 130 GeV. The MSSM and SM predictions for $\Gamma(H \longrightarrow \gamma\gamma)$ can differ by up to 10 % even in the decoupling limit of large m_A. We summarize recent progress in both the theoretical understanding of the background process $\gamma\gamma \longrightarrow q\overline{q}$, $q = \{b, c\}$, and in the expected detector performance allow for a high accuracy of the lightest MSSM or SM Higgs boson decay into a $b\overline{b}$ pair. We find that for optimized but still realistic detector and accelerator assumptions, statistically a 1.4% accuracy is feasible after about four years of collecting data for a Higgs boson mass which excludes $\tan\beta < 2$.

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1 Introduction

Uncovering the origin of particle masses is set to dominate the coming decade in high energy physics. It is commonly assumed that the Higgs mechanism [1] gives rise to all vector boson and fermion masses present in the Standard Model (SM) and indirect experimental evidence points to the existence of a light Higgs, possibly in the range predicted by the minimal supersymmetric SM (MSSM) [2]. The MSSM is attractive to theoreticians for several reasons, mostly, however, because supersymmetry stabilizes the quadratic divergences of scalars when the theory is extrapolated to GUT-scale energies. If one believes in such vast extrapolations, the SM-Higgs must be in a window between roughly 130-180 GeV due to restrictions on vacuum stability and the perturbative framework respectively [3, 4], not considering fine tuning (hierarchy) problems. Alternatively, a lighter Higgs mass would indicate the scale at which to expect new physics [4].

Supersymmetry and the requirement of anomaly cancellation both require al least two complex Higgs doublets in the MSSM leading to five physical degrees of freedom, two neutral CP-even (h,H), one neutral CP-odd (A) and two charged Higgs bosons ($H^{\pm}$) [2]. In contrast to a general two doublet Higgs model (2DHM), the MSSM Higgs sector has (at tree level) only two free parameters, commonly chosen as the ratio of the vacuum expectation values of up-and down-type Higgs bosons, $\tan\beta$, and the mass of the pseudoscalar Higgs, m_A. The lightest neutral scalar, h, in this model must be below 130 GeV [5].

Any supersymmetric extension of the SM must, in order to be phenomenologically viable, contain terms in the Lagrangian which break it. Commonly one introduces so called soft SUSY-breaking terms which can be thought to originate from supergravity or gauge mediation for instance [6]. In typical 'sugra' scenarios, the squark and gaugino scales are each degenerate at the SUSY-GUT scale and mass terms then evolve down to electroweak (EW) energies. EW-symmetry breaking in these models is then brought about by the Higgs boson mass parameter m_2 developing a negative vacuum expectation value [7].

While *a priori* a multitude of possibilities exists considering all the undetermined parameters of models beyond the SM, plausible SUSY scenarios often predict heavy pseudoscalar masses, $m_A > 400$ GeV [7]. At the LHC proton proton collider the sensitivity to these values of m_A is very limited. A negative Higgs search at LEP II would lead to a lower bound $\tan\beta > 2$ in the MSSM. Values of $\tan\beta > 3$ imply that for pseudoscalar masses above 250 GeV the LHC would not be able to provide information on m_A for intermediate values of $\tan\beta$ and above 500 GeV, only for very large $\tan\beta$ the $\tau^+\tau^-$ decay could be utilized [8].

In this context it is therefore of considerable interest to study observables which possess a large enough sensitivity to m_A in the decoupling limit compared to the SM prediction. The partial Higgs width $\Gamma(H \longrightarrow \gamma\gamma)$, measured at the $\gamma\gamma$ Compton-
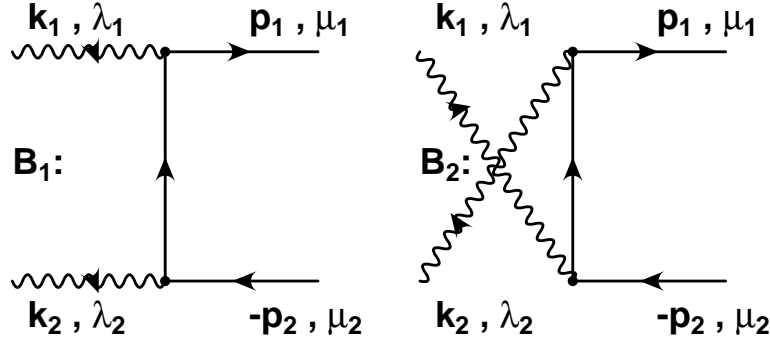
Figure 1: The Born process in polarized $\gamma + \gamma(J_z = 0) \longrightarrow q + \overline{q}$. QCD radiative corrections lead to large non-Sudakov double logarithms which need to be resummed.

backscattered option of a future linear $e^{\pm}$ collider, is such a quantity [14]. The decoupling limit of the diphoton Higgs width was studied in Ref. [15] and it was found that the MSSM and SM predictions can differ by up to 10 % in this limit. It was recently demonstrated in Ref. [16, 17] that using conservative assumptions an accuracy of 2% is feasible at a PLC. In the next section we briefly review the status of radiative corrections to both the signal (S) $\gamma\gamma \longrightarrow H \longrightarrow b\overline{b}$ and background (BG) process $\gamma\gamma \longrightarrow q\overline{q}$, $q = \{b, c\}$ and then present new Monte Carlo results assuming a slightly more optimistic (but still realistic) detector performance as in [16]. The results are summarized in section 4 in the above context of achievable precision measurements at a future photon photon collider.

2   Radiative Corrections

There has been considerable progress in the theoretical understanding of the BG to the intermediate mass Higgs boson decay into $b\overline{b}$ recently. The Born cross section in Fig. 1 for the $J_z = 0$ channel is suppressed by $\frac{m_q^2}{s}$ relative to the $J_z = \pm 2$ which means that by ensuring a high degree of polarization of the incident photons[1] one can *simultaneously* enhance the signal and suppress the background. QCD radiative corrections can remove this suppression, however, and large bremsstrahlung and double logarithmic corrections need to be taken into account.

In Ref. [10] the exact one loop corrections to $\gamma\gamma \longrightarrow q\overline{q}$ were calculated and the largest virtual correction was contained in novel non-Sudakov double logarithms. For some

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[1]Ratios of $J_0/J_2 = 20..50$ are feasible in presently considered designs, e.g. Ref. [18].
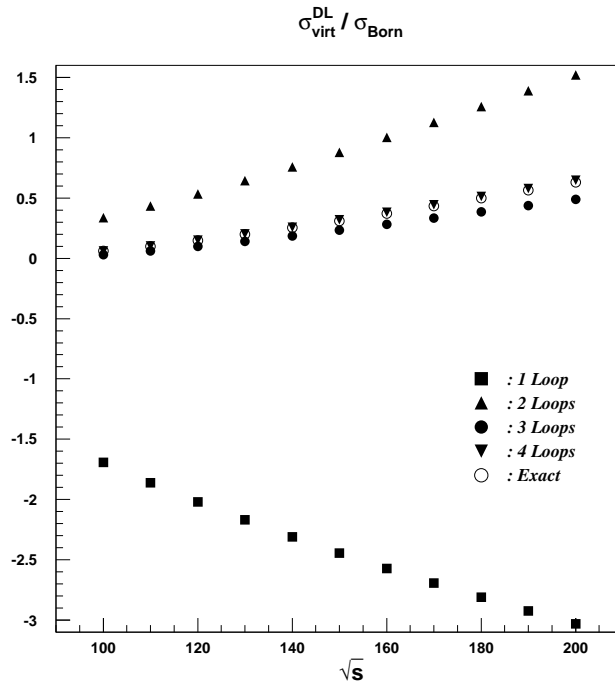
Figure 2: The size of the virtual non-Sudakov double logarithmic (DL) contribution relative to the Born cross section through four loops. The exact DL result (open circles) is given by the all orders resummation according to Ref. [20] and is in very good agreement with the four loop approximation given in Ref. [21]. The huge one and loop contributions can be seen to lead to physically distorted results.

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choices of the invariant mass cutoff y_{cut} even a negative cross section was obtained in this approximation. The authors of Ref. [19] elucidated the physical nature of the novel double logarithms and performed a two loop calculation in the DL-approximation. The results restored positivity to the physical cross section. In Ref. [20], three loop DL-results were presented which revealed a factorization of Sudakov and non-Sudakov DL's and led to the all orders resummation of all DL in form of a confluent hypergeometric function $_2F_2$. The general form of the expression is $\sigma_{DL} = \sigma_{Born}(1 + \mathcal{F}_{DL})\exp(\mathcal{F}_{Sud})$. Fig 2 demonstrates that at least four loops on the cross section level are required to achieve a converged DL result [21].

At this point the scale of the QCD-coupling is still unrestrained and differs by more than a factor of two in-between the physical scales of the problem, m_q and m_H. This uncertainty was removed in Ref. [22] by introducing a running coupling $\alpha_s(\mathbf{l}_\perp^2)$ into each loop integration (see Fig. 3), where $l_\perp$ denotes the perpendicular Sudakov loop momentum.

The effect of the RG-improvement lead to $\sigma_{DL}^{RG} = \sigma_{Born}(1 + \mathcal{F}_{DL}^{RG})\exp(\mathcal{F}_{Sud}^{RG})$ and the results are depicted in Fig. 4 for two choices of the gluon energy cutoff $l_c \equiv \epsilon\sqrt{s}$ compared to the theoretically allowed upper and lower limits of the DL-approximation evaluated at m_b^2 and m_H^2 respectively. In each case the RG-improved result remains inside the two DL-limits. The effective scale, defined simply as the one used in the DL-approximation which gives a result close to the RG-improved values, depends on ϵ, however in general is rather much closer to m_q than m_H [22].

On the signal side, the relevant radiative corrections have long been known up to NNL order in the SM [23, 24] and are summarized including the MSSM predictions in Ref. [8]. For our purposes the one loop corrections to the diphoton partial width depicted in Fig. 5 are sufficient as the QCD corrections are small in the SM. The important point to make here and also the novel feature in this analysis is that the branching ratio BR $(H \longrightarrow b\overline{b})$ is corrected by the same RG-improved resummed QCD Sudakov form as the continuum heavy quark background [16]. This is necessary in order to employ the same two jet definition for the final state. Since we use the renormalization group improved massive Sudakov form factor $\mathcal{F}_{Sud}^{RG}$ of Ref. [22], we prefer the Sterman-Weinberg jet definition [25] schematically depicted in Fig. 6. We also use an all orders resummed running quark mass evaluated at the Higgs mass for $\Gamma(H \longrightarrow b\overline{b})$. For the total Higgs width, we include the partial Higgs to $b\overline{b}, c\overline{c}, \tau^+\tau^-, WW^*, ZZ^*$ and gg decay widths with all relevant radiative corrections.
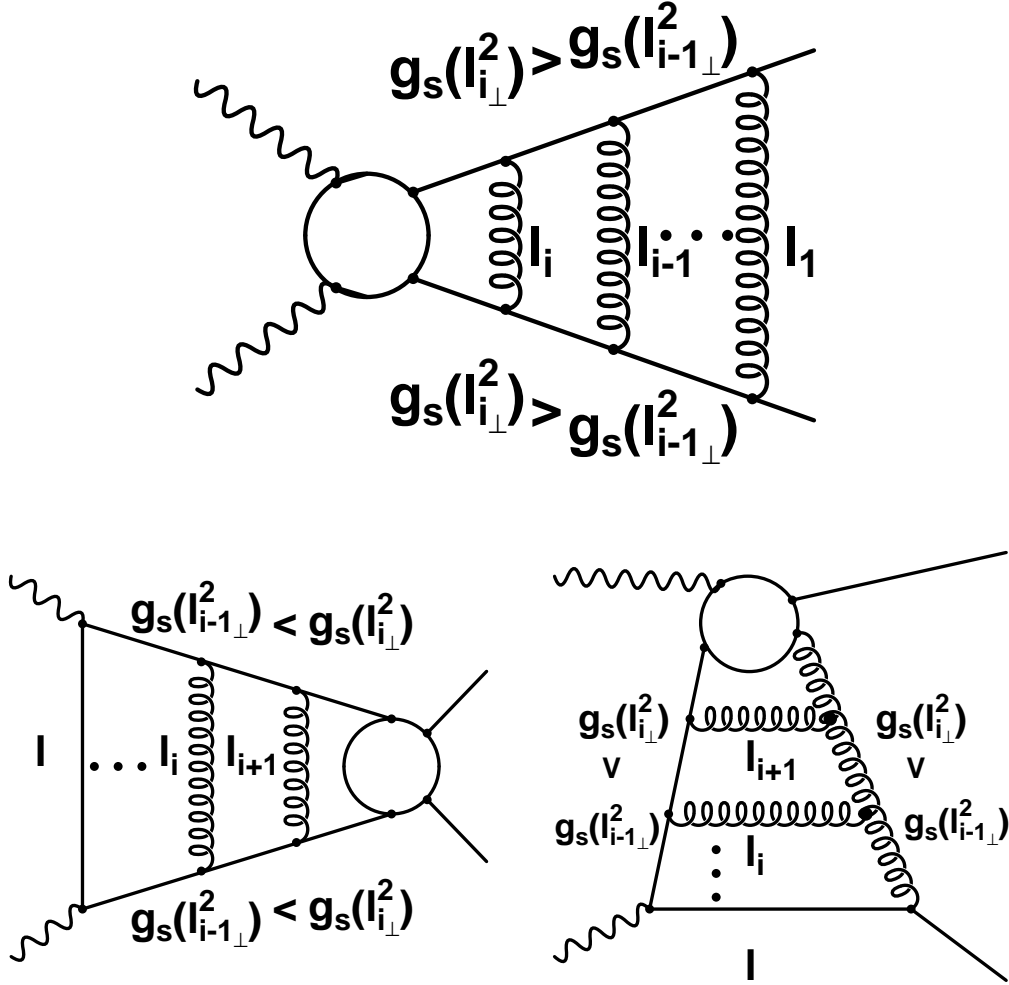
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Figure 3: The upper plot shows a schematic Feynman diagram leading to the Sudakov double logarithms in the process $\gamma\gamma(J_z = 0) \longrightarrow q\bar{q}$ with i gluon insertions. The blob denotes a hard momentum going through the omitted propagator in the DL-phase space. Crossed diagrams lead to a different ordering of the Sudakov variables with all resulting C_A terms canceling the DL-contributions from three gluon insertions [20]. The scale of the coupling $\alpha_s = \frac{g_s^2}{4\pi}$ is indicated at the vertices and explicitly taken into account in this work. The lower row depicts schematic Feynman diagrams leading to the renormalization group improved hard (non-Sudakov) double logarithms in that process. The topology on the left-hand diagram is Abelian like, and the one on the right is non-Abelian beyond one loop.
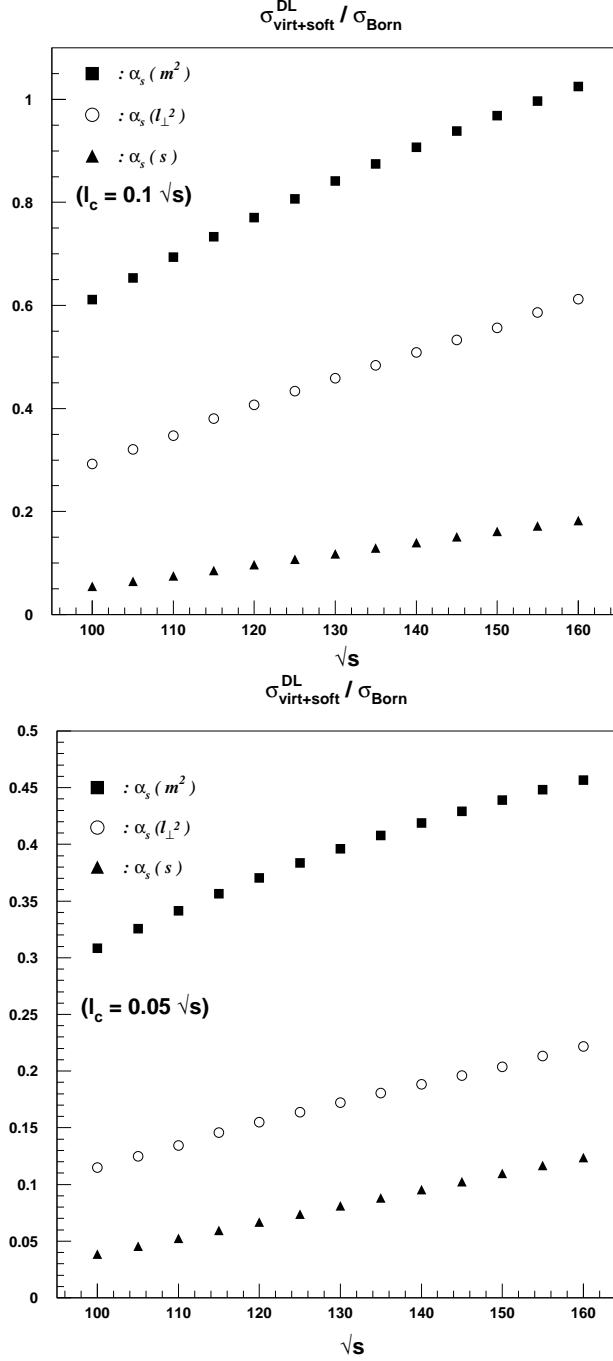
Figure 4: The effect of the renormalization group improved form factor (circles) of Fig. 3 in comparison to using the DL form factors with the indicated values of the strong coupling. The upper plot corresponds to $l_c = 0.1\sqrt{s}$ and the lower one to $l_c = 0.05\sqrt{s}$. The effect is displayed for the bottom quark with $m_b = 4.5$ GeV.
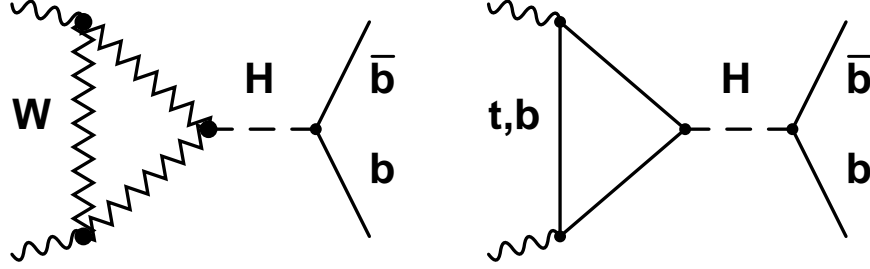
Figure 5: The Standard Model process $\gamma\gamma \longrightarrow H \longrightarrow b\bar{b}$ is mediated by $W-$boson and $t-$ and $b-$quark loops.
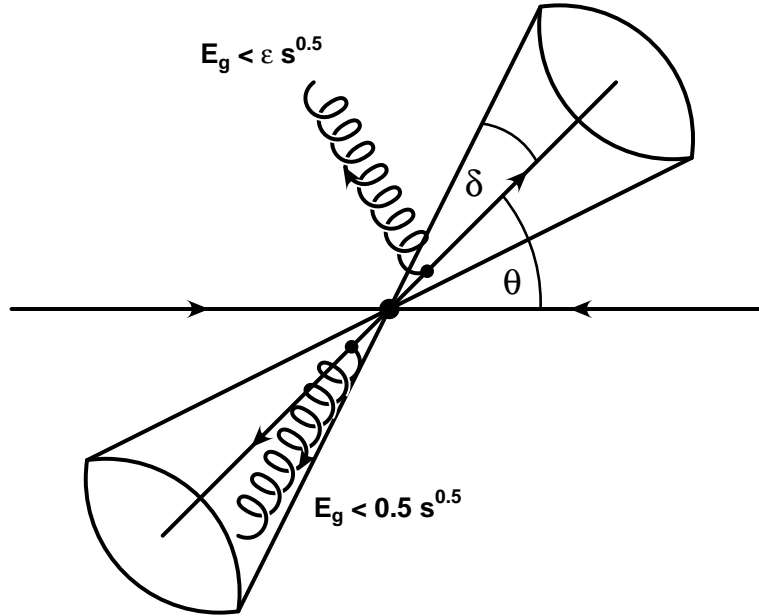

Figure 6: The parameters of the Sterman-Weinberg two-jet definition used in this work. Inside an angular cone of size δ arbitrary hard gluon bremsstrahlung is included. Radiation outside this cone is only permitted if the gluon energy is below a certain fraction (ϵ) of the incident center of mass energy. The thrust angle is denoted by θ.

3 Numerical Results

We begin with a few generic remarks concerning the uncertainties in our predictions. The signal process $\gamma\gamma \longrightarrow H \longrightarrow b\overline{b}$ is well understood and NNL calculations are available. The theoretical error is thus negligible [8].

There are two contributions to the background process $\gamma\gamma \longrightarrow q\overline{q}$ which we neglect in this paper. Firstly, the so-called resolved photon contribution [9] was found to be a small effect, e.g. [10], especially since we want to reconstruct the Higgs mass from the final two-jet measurements and impose angular cuts in the forward region. In addition the good charm suppression also helps to suppress the resolved photon effects as they give the largest contribution. The second contribution we do not consider here results from the final state configuration where a soft quark is propagating down the beam pipe and the gluon and remaining quark form two hard back-to back-jets [11]. We neglect this contribution here due to the expected excellent double b-tagging efficiency and the strong restrictions on the allowed acollinearity discussed below[2]. A good measure of the remaining theoretical uncertainty in the continuum background is given by scanning it below and above the Higgs resonance. For precision extractions of $\Gamma(H \longrightarrow \gamma\gamma)$ the exact functional form for resonant energies is still required, though.

In terms of possible systematic errors, the most obvious effect comes from the theoretical uncertainty in the bottom mass determination. Recent QCD-sum rule analyses, however, reach below the 2% level for $\overline{m}_b(\overline{m}_b)$ [12, 13]. For quantitative estimates of exptected systematic experimental errors it is clearly too early to speculate at this point. The philosophy adopted henceforth is that we assume that they can be neglected at the 1% level and concentrate purely on the statistical error.

We focus here not on specific predictions for cross sections, but instead on the expected statistical accuracy of the intermediate mass Higgs signal at a PLC. As detailed in Refs. [14], due to the narrow Higgs width, the signal event rate is proportional to $N_S \sim \left.\frac{dL_{\gamma\gamma}}{dw}\right|_{m_H}$, while the BG is proportional to $L_{\gamma\gamma}$.

To quantify this, we take the design parameters of the proposed TESLA linear collider [18, 26], which correspond to an integrated peak $\gamma\gamma$-luminosity of 15 fb^{-1} for the low energy running of the Compton collider. The polarizations of the incident electron beams and the laser photons are chosen such that the product of the helicities $\lambda_e\lambda_\gamma = -1$ [3]. This ensures high monochromaticity and polarization of the photon beams [18, 26]. Within this scenario a typical resolution of the Higgs mass is about 10 GeV, so that for comparison with the background process $BG \equiv \gamma\gamma \longrightarrow q\overline{q}$ one can use [14] $\frac{L_{\gamma\gamma}}{10\,\text{GeV}} = \left.\frac{dL_{\gamma\gamma}}{dw}\right|_{m_H}$ with $\left.\frac{dL_{\gamma\gamma}}{dw}\right|_{m_H} = 0.5$ fb^{-1}/GeV. The number of background events is then given by $N_{BG} = L_{\gamma\gamma}\sigma_{BG}$.

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[2]As discussed in Ref. [11] the B-hadrons from the slow b-quark could be dragged towards the gluon side and thus give rise to displaced decay vertices in the gluon jet. It may be of inter est to perform further systematic MC studies of this effect.

[3]The maximal initial electron polarization for existing projects is 85 %, e.g. Ref. [18].

In principle it is possible to use the exact Compton profile of the backscattered photons to obtain the full luminosity distributions. The number of expected events is then given as a convolution of the energy dependent luminosity and the cross sections. Our approach described above corresponds to an effective description of these convolutions, since these functions are not precisely known at present. Note that the functional forms currently used generally assume that only one scattering takes place for each photon, which may not be realistic. Once the exact luminosity functions are experimentally determined it is of course trivial to incorporate them into a Monte Carlo program containing the physics described in this paper.

In Ref. [16] it was demonstrated that in order to achieve a large enough data sample, a central thrust angle cut $|\cos\theta| < 0.7$ is advantageous and is adopted here. We also assume a (realistic) 70% double b-tagging efficiency. For the charm rejection rate, however, it seems now possible to assume an even better detector performance. The improvement comes from assuming a better single point resolution, thinner detector modules and moving the vertex detectors closer to the beam-line [27].

With these results in hand we keep $|\cos\theta| < 0.7$ fixed and furthermore assume the $c\bar{c}$ misidentification rate of 0.5%, (half of that in Ref. [16]). We vary the cone angle δ between narrow (10°), medium (20°) and large (30°) cone sizes for both $\epsilon = 0.1$ and $\epsilon = 0.05$. The upper row of Fig. 7 demonstrates that for the former choice of the energy cutoff parameter we achieve the highest statistical accuracy for the large $\delta = 30^o$ scenario of around 2%. We emphasize, however, that in this case also the missing $\mathcal{O}(\alpha_s^2)$ bremsstrahlung corrections could become important.

The largest effect is obtained by effectively suppressing the background radiative events with the smaller energy cutoff of $\epsilon = 0.05$ outside the cone (the inside is of course independent of ϵ). Here the lower row of Fig. 7 demonstrates that the statistical accuracy of the Higgs boson with $m_H < 130$ GeV can be below the 2% level after collecting one year of data. We should mention again that for this choice of ϵ we might have slightly enhanced the higher order (uncanceled) cutoff dependence. The dependence on the photon-photon polarization degree is visible but not crucial. Comparing with the results of Ref. [16] we also conclude that the new optimized charm misidentification rate leads to only slight improvements for $\sqrt{N_{tot}}/N_S$.

In summary, it seems very reasonable to expect that at the Compton collider option we can achieve a 2% statistical accuracy of an intermediate mass Higgs boson signal after collecting data over one year of running.

4   Conclusions

In this paper we have studied the Higgs signal and continuum background contributions to the process $\gamma\gamma \longrightarrow b\bar{b}$ at a high-energy Compton collider. We have used all relevant
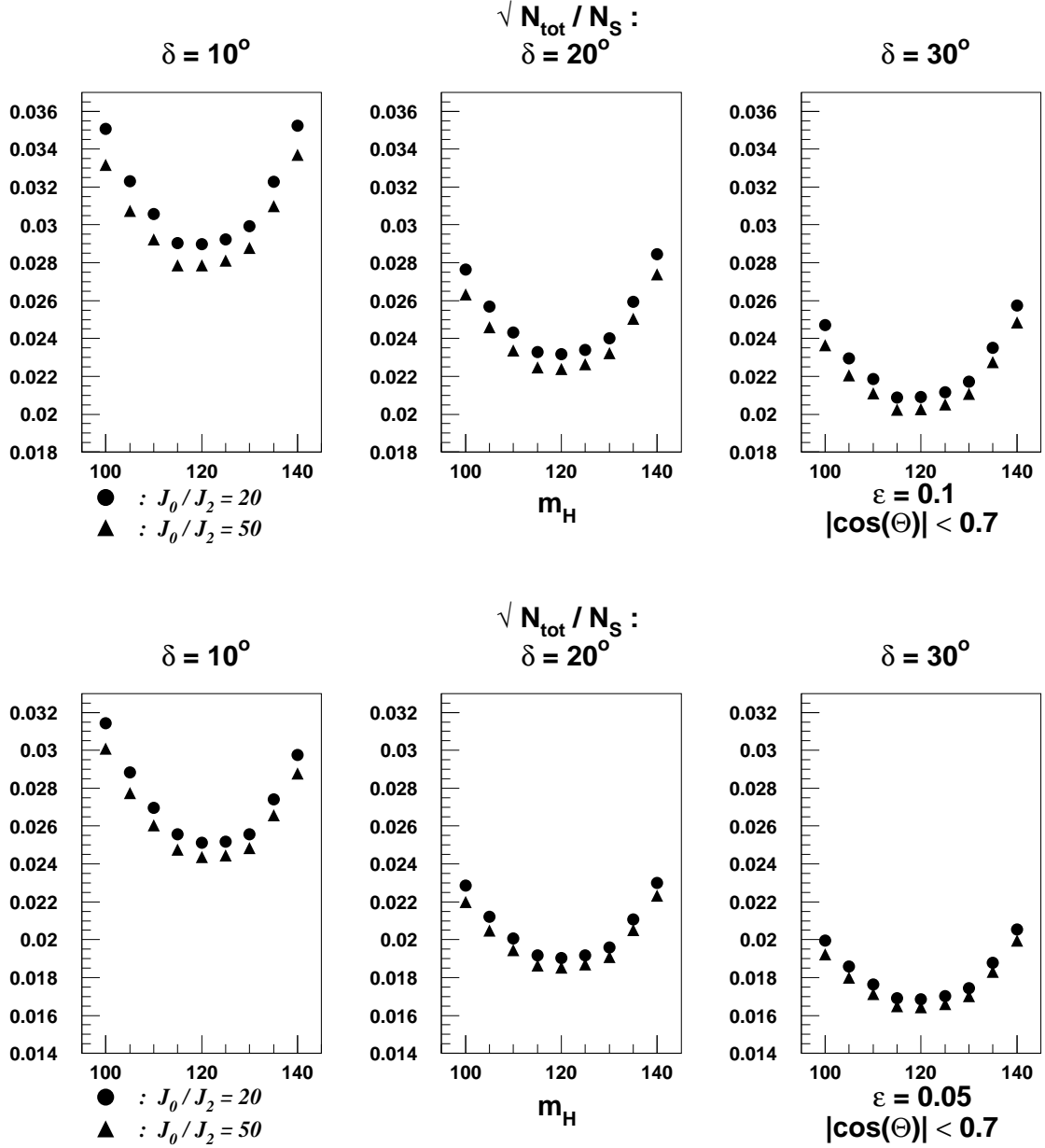
Figure 7: The cone-angle dependence of the inverse statistical significance of the intermediate mass Higgs signal for the displayed values of thrust and energy cut parameters. Overall a 70% double b-tagging efficiency and a 0.5% charm misidentification rate are assumed. For larger values of δ the number of events is enlarged, however, the theoretical uncertainty increases . For smaller values of ϵ higher order cutoff dependent terms might becom e important.

QCD radiative corrections to both the signal and BG production available in the literature. The Monte Carlo results using a variety of jet-parameter variations revealed that the intermediate mass Higgs signal can be expected to be studied with a statistical uncertainty between (excluding the narrow cone angle scenario) 2.4% in a realistic and 1.6% in an optimistic scenario after one year of collecting data for a Higgs-mass which excludes $\tan \beta < 2$.

Together with the expected uncertainty of 1% from the e^+e^- mode determination of BR($H \longrightarrow b\bar{b}$), and assuming *four* years of collecting data, we conclude that a measurement of the partial width $\Gamma(H \longrightarrow \gamma\gamma)$ of 1.4% precision level[4] is feasible for the MSSM mass range from a purely statistical point of view. This level of accuracy could significantly enhance the kinematical reach of the MSSM parameter space in the large pseudoscalar mass limit and thus open up a window for physics beyond the Standard Model.

For the total Higgs width, the main uncertainty is given by the error in the branching ratio BR($H \longrightarrow \gamma\gamma$), which at present is estimated at the 15 % level [28]. For Higgs masses above 110 GeV, the total Higgs width could be determined more precisely through the Higgs-strahlung process [29, 30] and its decay into WW^* [31]. This is only possible, however, if the supersymmetric lightest Higgs boson coupling to vector bosons is universal (i.e. the same for hWW and hZZ) and provided the optimistic luminosity assumptions can be reached.

In summary, using realistic and optimized machine and detector design parameters, we conclude that the Compton collider option at a future linear collider can considerably extend our ability to discriminate between the SM and MSSM scenarios.

Acknowledgments

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$g_s(I^2_{i-1_\perp}) < g_s(I^2_{i_\perp})$

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$g_s(I^2_{i-1_\perp}) < g_s(I^2_{i_\perp})$